STRANGENESS IN THE NUCLEON: THE STRANGE VECTOR FORM FACTORS

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ABSTRACT

We discuss two descriptions of the nucleon's strange vector form factors in the framework of vector meson dominance. The first, an updated and extended version of Jaffe's dispersion analysis, approximates the spectral functions of the form factors as a sum of vector meson poles, whereas the second combines vector meson dominance in the ω and ϕ meson sector with an intrinsic strangeness distribution from a kaon cloud.

KEYWORDS

Nucleon properties; strangeness; strange quark; dispersion theory; vector meson dominance.

STRANGENESS IN THE NUCLEON

The strangeness content of the nucleon, as probed by the matrix elements $\langle N|\overline{s}\Gamma s|N\rangle$ of strange quark operators $s, \bar{s}$ in a Lorentz channel specified by Γ, offers a key to intriguing and little understood quantum effects in the nucleon wave function. Due to the relatively small strange quark mass these effects can be sizable, and growing experimental evidence indeed indicates unexpectedly large amounts of strangeness in the nucleon. From pion-nucleon scattering data, for example, one can extract the ratio

$$R_s = \frac{\langle p|\overline{s}s|p\rangle}{\langle p|\overline{u}u + \overline{d}d + \overline{s}s|p\rangle} \tag{1}$$

(u, d and s are the up, down and strange quark fields, and $|p\rangle$ denotes the nucleon state), and obtains surprisingly large (although somewhat controversial) values, up to $R_s \simeq 0.1 - 0.2$ (Cheng

and Dashen, 1971; Cheng, 1976; Donoghue and Nappi, 1986; Gasser, Leutwyler and Sainio, 1991; Kluge, 1995). This implies that $\langle p|\overline{s}s|p\rangle$ can reach almost half the magnitude of the corresponding up-quark matrix element and that the nucleon mass would be reduced by $\approx 300\,\mathrm{MeV}$ in a world with massless strange quarks.

Further and more direct evidence for sizeable strange quark effects in the nucleon has emerged since the end of the eighties from deep inelastic μ-p scattering data. The European Muon Collaboration (EMC) measured the polarized proton structure function $g_1^p(x)$ in a large range of the Bjorken variable, $x \in [0.01, 0.7]$ (Ashman *et al.*, 1988, 1989) and found, after Regge extrapolation to $x = 0$ and combination with earlier SLAC data,

$$\int_0^1 dx\, g_1^p(x) = 0.126 \pm 0.010\,(stat) \pm 0.015\,(syst) \tag{2}$$

at $Q^2 = 10\,\mathrm{GeV}^2/c^2$. Without strange quark contributions one would expect, following Ellis and Jaffe (1974), a significantly larger value, 0.175 ± 0.018. The data therefore indicate a nonvanishing strange quark contribution $\Delta s = -0.16 \pm 0.008$ to the proton spin (if $SU(3)$ is not too badly broken), or equivalently, via the Bjorken sum rule, a substantial strangeness contribution $\langle p|\overline{s}\gamma_\mu\gamma_5 s|p\rangle$ to the proton matrix element of the isoscalar axial-vector current. The low-energy elastic ν–p scattering experiment E734 at Brookhaven (Ahrens *et al.*, 1987) complemented the EMC data by measuring the same matrix element at smaller momenta ($0.4\,\mathrm{GeV}^2 < Q^2 < 1.1\,\mathrm{GeV}^2$). The extracted axial vector current form factors are consistent with the muon scattering data.

The above experimental findings indicate a role of strange quarks in the nucleon that goes beyond naive quark model expectations (for a more complete discussion see (Alberg, 1995)) and have triggered further theoretical and experimental investigations. In view of the expected channel-dependence of the strange quark matrix elements (see, for example, (Ioffe and Karliner, 1990)), an important part of this activity is directed towards new channels and, in particular, to the vector channel. The vector current matrix element describes the nucleon's strangeness charge and current distributions (in analogy to the electromagnetic case) by Dirac and Pauli form factors,

$$\langle p'|\overline{s}\gamma_\mu s|p\rangle = \overline{N}(p')\left(\gamma_\mu F_1^{(s)}(q^2) + i\frac{\sigma_{\mu\nu}q^\nu}{2M_N}F_2^{(s)}(q^2)\right)N(p). \tag{3}$$

(N is the free Dirac spinor of the nucleon and $q = p' - p$.) Particularly attractive features of the vector matrix element are its scale independence (due to strangeness conservation, i.e. up to weak corrections) and its direct experimental accessibility via parity-violating lepton scattering off different hadronic targets. Several experiments of this type are in preparation at CEBAF and MAMI (Musolf *et al.*, 1994), and SAMPLE at Bates (McKeown and Beck, 1989) already started to take data. These experiments will, in fact, provide the first direct low-energy measurements of sea quark effects in hadrons.

The expected data will determine, in particular, the leading nonvanishing moments of the vector strangeness distribution, namely the strangeness radius r_s^2 and magnetic moment μ_s,

$$\mu_s = F_2^{(s)}(0) = G_M^{(s)}(0), \qquad r_s^2 = 6\frac{d}{dq^2}F_1^{(s)}(q^2)|_{q^2=0}, \qquad (r_s^2)_{Sachs} = 6\frac{d}{dq^2}G_E^{(s)}(q^2)|_{q^2=0}. \tag{4}$$

Note the two alternative definitions of the moments, which are both currently in use. One of them

is based on the Sachs form factors

$$G_E^{(s)}(q^2) = F_1^{(s)}(q^2) + \frac{q^2}{4M_N^2} F_2^{(s)}(q^2), \qquad G_M^{(s)}(q^2) = F_1^{(s)}(q^2) + F_2^{(s)}(q^2). \tag{5}$$

Since sea quark distributions in hadrons arise from a subtle interplay of quantum effects in QCD, their reproduction in hadron models is much more challenging than the calculation of the standard static observables. Reflecting these difficulties, present model calculations of the strange form factors (for a comparison see (Forkel *et al.*, 1994)) contain large and often uncontrolled theoretical uncertainties and are partially inconsistent with each other. Lattice calculations of strange sea quark distributions, on the other hand, are computationally very demanding and have not yet been carried out (see, however, (Liu and Dong, 1994)). The particular value of the strange quark mass, which is neither light nor heavy compared to the QCD scale Λ_{QCD}, further complicates the theoretical situation. In contrast to the light up and down quarks, the effects of the heavier strange quark are much harder to approach from the chiral limit, *i.e.* by an expansion in the quark mass. On the other hand, the strange quark is too light for the methods of the heavy-quark sector, *e.g.* the nonrelativistic approximation or the heavy-quark symmetry, to work.

In the following we discuss two theoretical approaches to the strange form factors which bypass, to a different degree, the need for detailed nucleon model calculations. They are both based on the phenomenologically successful vector meson dominance (VMD) concept, which is implemented in the following section in a dispersion theoretical framework (Jaffe, 1989; Forkel, 1995), and in the subsequent section via current field identities (Forkel *et al.*, 1994).

DISPERSION ANALYSIS

The dispersive approach, initiated by Jaffe (1989), permits a nucleon-model independent estimate of the strange form factors on the basis of phenomenological input. It starts from the dispersion relations

$$F_i^{(s)}(q^2) = \frac{1}{\pi} \int_{s_0}^{\infty} ds \, \frac{Im\{F_i^{(s)}(s)\}}{s - q^2}, \tag{6}$$

where $s_0 = (3m_\pi)^2$ is the three-pion threshold and subtraction terms are suppressed. The spectral functions $\pi^{-1} Im\{F_i^{(s)}\}$ receive contributions from intermediate (on-shell) states with $I^G J^{PC} = 0^- 1^{--}$, through which the strangeness current couples to the nucleon. In the pole approximation they are represented by N sharp vector meson states,

$$\frac{1}{\pi} Im\{F_i^{(s)}(s)\} = \sum_{v=1}^{N} B_i^{(v)} m_v^2 \, \delta(s - m_v^2), \qquad i \in \{1, 2\}. \tag{7}$$

This ansatz is perfectly adequate for the two lowest-lying, narrow-width resonances ω and ϕ. The additional poles are effectively summarizing strength from higher lying, broader resonances and from continuum contributions. Eq. (7) contains $3N$ a priori unknown mass and coupling parameters. Jaffe (1989) realized that the three lowest-lying masses and the couplings of the ω and ϕ poles can be estimated model-independently, since (i) there exists another current, the isoscalar electromagnetic current $J_\mu^{(I=0)}$, which carries the same quantum numbers as the strange

current and thus couples through the same intermediate states, (ii) the associated isoscalar form factors are well measured in a large range of q^2 and well fitted by a dispersive 3-pole ansatz (Höhler *et al.*, 1976; Mergell *et al.*, 1995) and (iii) the flavor structure of the first two poles is known.

The masses $m_1 - m_3$ in (7) can thus be identified with the pole positions found in the 3-pole fits to the electromagnetic form factors (in particular, $m_1 = m_\omega$ and $m_2 = m_\phi$). Furthermore, the four couplings $g(V, J)$ ($V = \omega, \phi$; $J = J^{(I=0)}$, $J^{(s)} \equiv \bar{s}\gamma s$) of the vector meson states

$$
\begin{aligned}
|\omega\rangle &= \cos\epsilon \, \frac{1}{\sqrt{2}} \left(|\overline{u}\gamma_\mu u\rangle + |\overline{d}\gamma_\mu d\rangle \right) - \sin\epsilon \, |\overline{s}\gamma_\mu s\rangle, \\
|\phi\rangle &= \sin\epsilon \, \frac{1}{\sqrt{2}} \left(|\overline{u}\gamma_\mu u\rangle + |\overline{d}\gamma_\mu d\rangle \right) + \cos\epsilon \, |\overline{s}\gamma_\mu s\rangle, \qquad (8)
\end{aligned}
$$

(the small angle $\epsilon = 0.053$ (Jain *et al.*, 1988) parametrizes the deviation from ideal mixing) to the neutral currents (defined via $\langle 0| \, J_\mu \, |V\rangle = g(V, J) \, m_V^2 \, \varepsilon_\mu$) are related by the assumption that the quark current of flavor i couples to the flavor-j component of the vector meson V with universal strength κ, and only for $i = j$, *i.e.*

$$
\langle 0| \, \overline{q}_i \gamma_\mu q_i \, | \, (\overline{q}_j q_j)_V \rangle = \kappa \, m_V^2 \, \delta_{ij} \, \varepsilon_\mu, \qquad (9)
$$

which works very well for the electromagnetic couplings. After parametrizing the vector-meson nucleon couplings as $g_i(\omega_0, N) = g_i \cos\eta_i$, $g_i(\phi_0, N) = g_i \sin\eta_i$ ($i = 1, 2$ denote the γ_μ and $\sigma_{\mu\nu} q^\nu$ couplings and ω_0, ϕ_0 the ideally mixed states) the four couplings $B_{1,2}^{(\omega,\phi)}$ in (7) can be obtained from the corresponding (fitted) isoscalar couplings $A_{1,2}^{(\omega,\phi)}$, which determine phenomenological values for η_i and κg_i:

$$
\begin{aligned}
A_i^{(\omega)} &= \tfrac{\kappa}{\sqrt{6}} \sin(\theta_0 + \epsilon) g_i \cos(\eta_i + \epsilon), \qquad & B_i^{(\omega)} &= -\kappa \sin\epsilon g_i \cos(\eta_i + \epsilon), \qquad (10) \\
A_i^{(\phi)} &= -\tfrac{\kappa}{\sqrt{6}} \cos(\theta_0 + \epsilon) g_i \sin(\eta_i + \epsilon), \qquad & B_i^{(\phi)} &= \kappa \cos\epsilon g_i \sin(\eta_i + \epsilon). \qquad (11)
\end{aligned}
$$

(θ_0 is the "magic angle" with $\sin^2\theta_0 = 1/3$.) Since the flavor content of the strength associated with m_3 is unknown, the above strategy cannot be applied to the couplings $B_i^{(3)}$. They are fixed instead by imposing weak constraints on the asymptotic behavior,

$$
\lim_{q^2 \to -\infty} q^{2(i-1)} F_i^s(q^2) \to 0 \qquad \Rightarrow \qquad \sum_v B_1^{(v)} = 0, \quad \sum_v B_2^{(v)} m_v^2 = 0, \qquad (12)
$$

which also normalize F_1. Jaffe's analysis included the minimal number of 3 poles in (7) and took the couplings $A_i^{(v)}$ from the almost twenty year old fits 8.1, 8.2 and 7.1 of Höhler *et al.* (1976) (with $m_3 = \{1.4, 1.8, 1.67\}$GeV, respectively). The resulting moments, averaged over the fits, were $r_s^2 = (0.16 \pm 0.06)\,\mathrm{fm}^2$, $(r_s^2)_{Sachs} = (0.14 \pm 0.07)\,\mathrm{fm}^2$ and $\mu_s = -(0.31 \pm 0.09)$. The indicated error estimates originate solely from the spread between the fits and are thus at best a rough lower bound on the complete error.

A new 3-pole fit to the current world data set of the electromagnetic form factors (Mergell *et al.*, 1995) has prompted our update (Forkel, 1995) of the strange form factor analysis. Besides being based on a considerably expanded data set, the fits of Mergell *et al.* are designed to reproduce the logarithmic QCD corrections to the form factor asymptotics, which partially originate from continuum contributions. Also, they find the third pole mass at the value of another well

established resonance in the isoscalar channel, $m_3 = 1.6\,\mathrm{GeV}$. The strange form factor analysis benefits from these additional features, since they increase the reliability of the extracted mass and coupling parameters. The updated values of the strangeness radius and magnetic moment are (Forkel, 1995)

$$r_s^2 = 0.22\,\mathrm{fm}^2, \qquad (r_s^2)_{Sachs} = 0.20\,\mathrm{fm}^2, \qquad \mu_s = -0.26. \tag{13}$$

While the square radius becomes 40 % larger than that found by Jaffe, $|\mu_s|$ is reduced by about 20 %. The bulk of the changes in r_s^2 and μ_s can be traced to differences in the ϕ-nucleon couplings of the used fits. Note that the estimates (13) are surprisingly large, of the order of the neutron charge radius $r_n^2 = -0.11\,\mathrm{fm}^2$ and the isoscalar magnetic moment of the nucleon, $\mu^{(I=0)} = 0.44$, respectively.

The momentum dependence of the 3-pole form factors at spacelike $Q^2 \equiv -q^2$ is shown in Figs. 1a and 1b. It reflects the sizes of the pole couplings: since the $|B_i^{(\omega)}|$ are about an order of magnitude smaller than the $|B_i^{(\phi)}|$, the leading $1/q^2$ dependence of the ϕ pole cannot be cancelled (as required by (12)) by the ω pole alone. Thus the coupling of the third pole must be of similar magnitude as the ϕ coupling, but of opposite sign. One therefore expects a dipole form of F_2^s with a mass parameter between m_2 and m_3, and an almost perfect fit for all space-like momenta is indeed obtained by the simple parametrization

$$F_1^{(s)}(q^2) = \frac{1}{6}\frac{r_s^2 q^2}{(1-\frac{q^2}{M_1^2})^2}, \qquad F_2^{(s)}(q^2) = \frac{\mu_s}{(1-\frac{q^2}{M_2^2})^2}, \tag{14}$$

with $M_1 = 1.31\,\mathrm{GeV} \simeq M_2 = 1.26\,\mathrm{GeV}$ (for Mergell *et al.*s parameters, i.e. with $r_s^2 = 5.680\,\mathrm{GeV}^{-2}$ and $\mu_s = -0.257$), which explicitly realizes the asymptotic behavior (12).

The 3-pole ansatz cannot, however, reproduce the asymptotic power behavior of the form factors established via QCD dimensional counting rules[1] (Brodsky and Farrar, 1975; Lepage and Brodsky, 1980). Ultimately, at very large, spacelike q^2, the form factors are dominated by extrinsic contributions, which originate from the renormalization of the strange current, are thus suppressed by higher powers of α, and decay as

$$F_1^{(s)}(q^2) \to \left(-q^2\right)^{-2}, \qquad F_2^{(s)}(q^2) \to \left(-q^2\right)^{-3}. \tag{15}$$

However, enforcing this behavior might not necessarily be the best choice for an optimal description of the form factors at *small and intermediate* momentum transfers in the pole approximation. The reason is that the other, intrinsic contributions, which originate from $s\bar{s}$ admixtures to the nucleon wave function, are, although asymptotically subleading,

$$F_1^{(s)}(q^2) \to \left(-q^2\right)^{-4}, \qquad F_2^{(s)}(q^2) \to \left(-q^2\right)^{-5}, \tag{16}$$

not α-suppressed. There might thus exist an intermediate range of momentum transfers where the form factors show the intrinsic decay behavior. Up to these momenta, the form factors would then be better described by enforcing the intrinsic behavior. Furthermore, the pole approximation is more reliable at smaller momenta, where also the deviation from the extrinsic behavior in the

[1] I am indebted to Stan Brodsky for helpful correspondence on this issue.

asymtotic tail would be of little effect. In the following we will briefly discuss the two minimal (4- and 6-pole) ansätze which can describe the extrinsic or intrinsic asymptotics (for more details see (Forkel, 1995)). In the framework of eq. (7) the extrinsic power behavior (15) requires minimally 4 poles,

$$F_i^{(s)}(q^2) = \sum_{v=1}^{4} B_i^{(v)} \frac{m_v^2}{m_v^2 - q^2}, \qquad i \in \{1,2\}, \tag{17}$$

together with the normalization and asymptotic constraints

$$\left(\begin{array}{cc} m_3^2 & m_4^2 \\ m_3^4 & m_4^4 \end{array}\right)\left(\begin{array}{cc} m_3^{-2} & 0 \\ 0 & m_4^{-2} \end{array}\right)\left(\begin{array}{c} B_1^{(3)} \\ B_1^{(4)} \end{array}\right) = - \left(\begin{array}{c} C_1^{(3)} \\ C_1^{(4)} \end{array}\right), \tag{18}$$

$$\left(\begin{array}{cc} m_3^2 & m_4^2 \\ m_3^4 & m_4^4 \end{array}\right)\left(\begin{array}{c} B_2^{(3)} \\ B_2^{(4)} \end{array}\right) = - \left(\begin{array}{c} C_2^{(3)} \\ C_2^{(4)} \end{array}\right), \tag{19}$$

($C_1^{(i)} \equiv \sum_{j=1}^{2} B_1^{(j)} m_j^{2(i-3)}$, $C_2^{(i)} \equiv \sum_{j=1}^{2} B_2^{(j)} m_j^{2(i-2)}$), which have a unique solution for the couplings $B_i^{(3,4)}$ as a function of the masses $m_{3,4}$. These solutions leave the value of the fourth mass, m_4, free. Maintaining the third pole at $m_3 = 1600\,\mathrm{MeV}$ and requiring m_4 to be larger than m_3 by at least a typical width of $\sim 300\,\mathrm{MeV}$ (so that it can be resolved in zero-width approximation), i.e. $m_4 \geq 1.9\,\mathrm{GeV}$, the results for the strangeness radius and magnetic moment interpolate smoothly and monotonically in the range

$$\begin{array}{rcl} 0.15\,\mathrm{fm}^2 \leq & r_s^2 & \leq 0.22\,\mathrm{fm}^2, \\ 0.14\,\mathrm{fm}^2 \leq & (r_s^2)_{Sachs} & \leq 0.20\,\mathrm{fm}^2, \\ -0.18 \geq & \mu_s & \geq -0.26. \end{array} \tag{20}$$

For $m_4 \to \infty$ the fourth pole does affect the momentum dependence only at $Q^2 >> m_4^2$, and for smaller Q^2 the 4-pole ansatz becomes identical to the 3-pole ansatz, which provides the upper bounds on r_s^2 and $|\mu_s|$ in (20). A fourth pole in the 2 GeV region, however, reduces the 3-pole moments by about a third. For all admissable values of m_4, the couplings of the third pole are necessarily large. In the 4-pole ansatz with $m_4 \sim 2\,\mathrm{GeV}$, also the coupling of the fourth pole is of comparable size and we expect quadrupole form factors. Indeed, the conservative choice $m_4 = 1.9\,\mathrm{GeV}$ is well fitted by $F_1^{(s)}(q^2) = (r_s^2 q^2/6)(1 - q^2/M_1^2)^{-3}$, $F_2^{(s)}(q^2) = \mu_s(1 - q^2/M_2^2)^{-3}$ with $r_s^2 = 0.1482\,\mathrm{fm}^2, \mu_s = -0.1789, M_1 = 1.47\,\mathrm{GeV}$ and $M_2 = 1.43\,\mathrm{GeV}$, which provide a lower bound on the (absolute) magnitude of the form factors with extrinsic asymptotics.

We now turn to the intrinsic asymptotics (16), which requires two more superconvergence relations for each form factor and minimally 6 poles, since the 5-pole ansatz would be overconstrained (Forkel, 1995). The relevant expressions are direct generalizations of (17), (18) and (19) from $N = 4$ to $N = 6$. As for the 4-pole ansatz, the couplings can be uniquely expressed in terms of the masses, leaving two more pole positions, m_5 and m_6, undetermined. Again we estimate the range of possible 6-pole form factors by requiring spacings of minimally 300 MeV between higher-lying poles. The most conservative estimate corresponds then to the mass values ($\{m_4, m_5, m_6\} = \{1.9, 2.2, 2.5\}\,\mathrm{GeV}$) and is again well fitted by the simplest one-mass-parameter formulae which match their asymptotic behavior: $F_1^{(s)}(q^2) = (r_s^2 q^2/6)(1 - q^2/M_1^2)^{-5}$, $F_2^{(s)}(q^2) = \mu_s(1 - q^2/M_2^2)^{-5}$, with $r_s^2 = 0.08879\,\mathrm{fm}^2, \mu_s = -0.1136, M_1 = 1.72\,\mathrm{GeV}$ (or somewhat better for $Q^2 \leq 10\,\mathrm{GeV}$ by $M_1 = 1.79\,\mathrm{GeV}$) and $M_2 = 1.66\,\mathrm{GeV}$. Increasing m_6, m_5 and m_4 up to infinity one arrives again

at the 3-pole form factors as an upper bound. The range of values for the leading moments is now given by

$$\begin{aligned} 0.089\,\mathrm{fm}^2 \leq \quad & r_s^2 & \leq 0.22\,\mathrm{fm}^2, \\ 0.081\,\mathrm{fm}^2 \leq \quad & (r_s^2)_{Sachs} & \leq 0.20\,\mathrm{fm}^2, \\ -0.086 \geq \quad & \mu_s & \geq -0.26. \end{aligned} \tag{21}$$

The intrinsic asymptotics can thus reduce the size of the strangeness radius and magnetic moment by about a factor of 3 relative to that of the 3-pole estimate. Comparing the different asymptotics we arrive at some general conclusions: (i) the r_s^2- and $|\mu_s|$–values of the minimal 3-pole ansatz are upper bounds and very likely overestimated, possibly by up to a factor of three, (ii) their values contain (at least in the realm of the pole approximation) significant information on the asymptotic behavior and its onset, and (iii) the pole approximation leads quite generally to *positive* strangeness radii and negative magnetic moments. The last point can be readily understood from the generic N–pole expressions

$$r_s^2 = \sum_v^N \frac{B_1^{(v)}}{m_v^2}, \qquad \mu_s = \sum_v^N B_2^{(v)}, \tag{22}$$

since the large ϕ couplings are positive in $F_1^{(s)}$ and negative in $F_2^{(s)}$ and since both the alternating signs of the couplings (due to the superconvergence relations) and the m_v^{-2} factor in r_s^2 suppress higher-pole contributions. The sign of the strangeness radius might, however, be changed by (*e.g.* $K\bar{K}$) continuum contributions, as discussed in the following sections.

GENERALIZED VECTOR MESON DOMINANCE

The pole approximation of the above dispersive analysis includes effects of the $\bar{K}K$ continuum (and those from the other cuts) at best implicitly. In the present section we present a model for the strange form factors which is based on a more general version of VMD and contains explicit contributions from the kaon cloud of the nucleon. This approach relies exclusively on the lightest, narrow isoscalar vector mesons, ω and ϕ, and is discussed in detail by Forkel *et al.* (1994). The VMD hypothesis is formulated in terms of current field identities (CFIs) (Kroll, Lee and Zumino, 1967), which imply the proportionality of the electromagnetic current to the field operators of the light, neutral vector mesons with the same quantum numbers. In the isocalar electromagnetic channel the CFI reads

$$J_\mu^{(I=0)} = A_\omega\, m_\omega^2\, \omega_\mu + A_\phi\, m_\phi^2\, \phi_\mu, \tag{23}$$

with the couplings A_ω, A_ϕ yet to be fixed. Generalizing VMD to the strangeness current, we write an analogous CFI

$$J_\mu^{(s)} = \bar{s}\gamma_\mu s = B_\omega\, m_\omega^2\, \omega_\mu + B_\phi\, m_\phi^2\, \phi_\mu. \tag{24}$$

After combining eqs. (23) and (24) into a vector equation, the couplings form the elements of a matrix $\hat{C}$. Sandwiching the CFIs between the physical vector meson states and the vacuum, and using eq. (9), we obtain an explicit form for $\hat{C}$,

$$\hat{C}_{I=0,s}(\epsilon) = \begin{pmatrix} A_\omega & A_\phi \\ B_\omega & B_\phi \end{pmatrix} = \kappa \begin{pmatrix} \frac{1}{\sqrt{6}}\sin(\theta_0+\epsilon) & \frac{-1}{\sqrt{6}}\cos(\theta_0+\epsilon) \\ -\sin\epsilon & \cos\epsilon \end{pmatrix}. \tag{25}$$

The CFIs lead to a general expression for the form factors. To derive it, we first note that eqs. (23) and (24), together with the requirement of strangeness and hypercharge conservation, imply $\partial^\mu V_\mu = 0$ (V_μ stands for either ω_μ or ϕ_μ), which simplifies the field equations to

$$(\Box + m_V^2)\, V_\mu = J_\mu^{(V)} \tag{26}$$

and therefore also implies that the vector meson source currents are conserved ($\partial^\mu J_\mu^{(V)} = 0$). We now take nucleon matrix elements of the field equations (26) and use the CFI's to write

$$\begin{pmatrix} \langle N(p')|\, J_\mu^{(I=0)}\, |N(p)\rangle \\ \langle N(p')|\, J_\mu^{s}\, |N(p)\rangle \end{pmatrix} = \hat{C}_{I=0,s}(\epsilon) \begin{pmatrix} \frac{m_\omega^2}{m_\omega^2-q^2} & 0 \\ 0 & \frac{m_\phi^2}{m_\phi^2-q^2} \end{pmatrix} \begin{pmatrix} \langle N(p')|\, J_\mu^{(\omega)}\, |N(p)\rangle \\ \langle N(p')|\, J_\mu^{(\phi)}\, |N(p)\rangle \end{pmatrix}. \tag{27}$$

Is is convenient to reexpress the vector meson source currents as linear combinations of currents with the same $SU(3)$ transformation behavior as $J^{(I=0)}$ and $J^{(s)}$, which we will denote as *intrinsic* (J_{in}). After furthermore separating the nucleon matrix elements into form factors, according to eq. (3) and its analog for $J_\mu^{(I=0)}$, we obtain our general VMD expression for the form factors:

$$\begin{pmatrix} F^{I=0}(q^2) \\ F^{(s)}(q^2) \end{pmatrix} = \hat{C}_{I=0,s}(\epsilon) \begin{pmatrix} \frac{m_\omega^2}{m_\omega^2-q^2} & 0 \\ 0 & \frac{m_\phi^2}{m_\phi^2-q^2} \end{pmatrix} \hat{C}^{-1}_{I=0,s}(\epsilon) \begin{pmatrix} F_{in}^{I=0}(q^2) \\ F_{in}^{(s)}(q^2) \end{pmatrix}. \tag{28}$$

According to their definition, the intrinsic form factors describe the extended source current distribution of the nucleon to which the vector mesons couple. Since both $J^{(I=0)}$, $J^{(s)}$ and their intrinsic counterparts $J_{in}^{(I=0)}$, $J_{in}^{(s)}$ are conserved, the full and the intrinsic form factors in eq. (28) have the same normalization at $q^2 = 0$. Combining eqs. (25) and (28), we finally obtain

$$\begin{pmatrix} F^{I=0}(q^2) \\ F^{(s)}(q^2) \end{pmatrix} = \begin{pmatrix} \frac{m_\omega^2}{m_\omega^2-q^2}\frac{\sin(\theta_0+\epsilon)\cos\epsilon}{\sin\theta_0} - \frac{m_\phi^2}{m_\phi^2-q^2}\frac{\cos(\theta_0+\epsilon)\sin\epsilon}{\sin\theta_0} & \frac{\cos(\theta_0+\epsilon)\sin(\theta_0+\epsilon)}{\sqrt{6}\sin\theta_0}\left(\frac{m_\omega^2}{m_\omega^2-q^2} - \frac{m_\phi^2}{m_\phi^2-q^2}\right) \\ \frac{\sqrt{6}\cos\epsilon\sin\epsilon}{\sin\theta_0}\left(\frac{m_\phi^2}{m_\phi^2-q^2} - \frac{m_\omega^2}{m_\omega^2-q^2}\right) & \frac{m_\phi^2}{m_\phi^2-q^2}\frac{\cos\epsilon\sin(\theta_0+\epsilon)}{\sin\theta_0} - \frac{m_\omega^2}{m_\omega^2-q^2}\frac{\sin\epsilon\cos(\theta_0+\epsilon)}{\sin\theta_0} \end{pmatrix}$$

$$\times \begin{pmatrix} F_{in}^{I=0}(q^2) \\ F_{in}^{(s)}(q^2) \end{pmatrix}. \tag{29}$$

Note that eq. (29) is independent of the overall vector-meson-current and vector-meson-nucleon coupling constants, which cancel each other due to charge normalization.

Up to now our discussion has been rather general, and different choices for the intrinsic form factors can be implemented in the given framework, as long as they do not lead to double counting with the VMD sector. Here, we adopt the kaon loop model of Musolf and Burkhardt (1993) for the intrinsic strangeness form factor (but use the physical value for the Λ mass instead of their flavor-symmetric value). This model describes the current-nucleon vertex corrections due to K–Λ loop graphs. Although the latter are U.V. finite, the loop momenta are cut off by meson-nucleon vertex form factors $H(k^2) = (m_K^2 - \Lambda^2)/(k^2 - \Lambda^2)$ from the Bonn potential (Holzenkamp *et al.*, 1989), since the effective hadronic description of the underlying physics breaks down at large momenta. The Bonn values for the cutoff Λ in the $N\Lambda K$ vertex were extracted from fits to baryon-baryon scattering data and lie in the range of $1.2 - 1.4$ GeV. One finds three amplitudes, $\Gamma_\mu^{(B,M,W)}$, which contribute to the intrinsic form factors. They are associated with processes where the current couples either to the baryon line (B), the meson line (M) or the meson-baryon vertex (V) in the

loop, and the intrinsic strange form factors are obtained from the nucleon matrix element of their sum,

$$\overline{N}(p')\left[\Gamma^B_\mu(p',p)+\Gamma^M_\mu(p',p)+\Gamma^V_\mu(p',p)\right]N(p)=\overline{N}(p')\left(\gamma_\mu F^{(s)}_{1,in}(q^2)+i\frac{\sigma_{\mu\nu}q^\nu}{2M_N}F^{(s)}_{2,in}(q^2)\right)N(p). \quad (30)$$

Explicit expressions for the $\Gamma^{(B,V,M)}$ are given in (Forkel *et al.*, 1994), where it is also shown that their sum satisfies the Ward-Takahashi identity. The values of the coupling and masses are fixed at $M_N = 939\,\mathrm{MeV}$, $M_\Lambda = 1116\,\mathrm{MeV}$, $m_K = 496\,\mathrm{MeV}$ and $g_{N\Lambda K}/\sqrt{4\pi} = -3.944$ (Holzenkamp *et al.*, 1989). Finally, we extract the *intrinsic isoscalar* form factors from the fit to the measured isoscalar form factors by inverting the VMD matrix in eq. (29). The full strangeness form factors are then determined by the second row of eq. (29). (The contribution from the intrinsic strangeness part to the isoscalar form factor is very small and plays almost no role in the determination of $F^{I=0}_{in}(q^2)$.)

In the generalized VMD approach the strangeness magnetic moment originates solely from the intrinsic kaon loop contribution (*cf.* eq. (28)), while both the Dirac and the Sachs strangeness radii get an additive contribution from the vector mesons. One finds

$$r_s^2 = -(0.0243 - 0.0245)\,\mathrm{fm}^2, \quad (r_s^2)_{Sachs} = -(0.040 - 0.045)\,\mathrm{fm}^2, \quad \mu_s = -(0.24 - 0.32). \quad (31)$$

The signs of the strangeness radii are the same as those of the intrinsic contribution.

The momentum dependence of the resulting Dirac and Pauli form factors is shown in Figs. 2a and 2b. For comparison, we also show the 4-pole and 6-pole form factors of Figs. 1a and 1b. The two approaches differ in sign and magnitude of the Dirac form factor (and thus of the strangeness radius), but yield more similar Pauli form factors. (The Pauli form factor from generalized VMD is almost identical to that of the 3-pole ansatz.) The combined data from SAMPLE (which measures $F_2^{(s)}$ at $Q^2 = 0.1\,\mathrm{GeV}^2$), CEBAF (in particular the G0 experiment (Beck *et al.*, 1991), which will measure G^s_M in the momentum range $0.1\,\mathrm{GeV}^2 < Q^2 < 0.5\,\mathrm{GeV}^2$ with a resolution $\delta\mu_s \simeq \pm 0.22$ at low Q^2) and MAMI (Heinen-Konschak *et al.*, 1993) should thus be sufficient to distinguish between the two VMD approaches. Note, finally, that the strangeness radius from generalized VMD is proportional to the sine of the mixing angle ϵ and would thus not receive any contribution from ideally mixed vector meson states, whereas in the dispersion analysis r_s^2 gets bigger as the mixing angle ϵ goes to zero, since the overall strangeness of the intrinsic charge distribution vanishes.

SUMMARY AND CONCLUSIONS

The vector meson dominance mechanism has a largely generic character and successfully describes electromagnetic interactions of hadrons. It should therefore be a useful starting point for estimates of the nucleon's strange vector form factors. We have discussed two such estimates: the first, a dispersive treatment, relies on phenomenological input and is nucleon-model independent, while the second builds on current field identities and consistently includes an intrinsic strangeness distribution due to the nucleon's kaon cloud.

The dispersive analysis, which we restrict to the pole approximation, is based on input from the isoscalar electromagnetic form factor data. We show that the minimal parametrization of the

spectral functions, Jaffe's 3-pole ansatz, yields upper bounds for the magnitude of the strangeness radius r_s^2 and magnetic moment μ_s. Using new fits to the current world data set of the isoscalar form factors, we also update the results of the 3-pole analysis and find that r_s^2 increases by 40 % while $|\mu_s|$ decreases by 20 %, compared to Jaffe's original values. Due to the unrealistic asymptotics of the 3-pole ansatz, however, these upper bounds do probably overestimate the moments: the reproduction of the asymptotic behavior derived from QCD counting rules requires additional poles and leads to significantly reduced values for the moments. Implementing the leading QCD asymptotics with a fourth pole term and a conservative estimate for the bulk position of higher-lying strength reduces the 3-pole results by more than a third, to $r_s^2 = 0.15\,\mathrm{fm}^2$, $(r_s^2)_{Sachs} = 0.14\,\mathrm{fm}^2$, $\mu_s = -0.18$, while the stronger decay due to intrinsic contributions leads to a further reduction by up to 50 %, $r_s^2 = 0.089\,\mathrm{fm}^2$, $(r_s^2)_{Sachs} = 0.081\,\mathrm{fm}^2$, $\mu_s = -0.086$.

Due to the alternating signs ("bump-dip" structure) of neighboring pole couplings, the momentum dependence of the pole-ansatz form factors turns out to be well fitted by simple multipole parametrizations. The positive sign of r_s^2 and the negative sign of μ_s follow the signs of the large ϕ couplings and are thus generic in the pole framework. The fact, furthermore, that the lightest relevant (i.e. K–Λ) intermediate states in the nucleon wave function produce the opposite, negative sign for r_s^2 may indicate that the K–$\bar{K}$ continuum could change the sign of the strangeness radius in the dispersive analysis.

This sign change indeed takes place in a generalization of the VMD framework which consistently includes kaonic contributions on the basis of current field identities. The latter permit to account explicitly for an intrinsic strangeness distribution of the nucleon, which we generate by a K–Λ loop amplitude. Restricting the VMD sector to the sharp ω and ϕ resonances, we find a negative and by half an order of magnitude smaller strangeness radius ($r_s^2 = -0.024\,\mathrm{fm}^2$, $(r_s^2)_{Sachs} \approx -0.043\,\mathrm{fm}^2$), while the strange magnetic moment predictions are of the same sign and a similar (loop-cutoff dependent) magnitude than those of the dispersion analysis, $\mu_s = -(0.24 - 0.32)$.

Both the dispersive and the CFI-based analysis allow many refinements. Perhaps the most important ones are the inclusion of $K\bar{K}$ continuum contributions in the former and a more complete model for the intrinsic form factors (Forkel, Musolf and Nielsen, 1996) in the latter. Investigations in both directions are in progress.

I would like to thank Professor Amand Faessler and the co-organizers of the school for creating such a stimulating, well-rounded and productive meeting. I would also like to thank Tom Cohen, Xuemin Jin and Marina Nielsen for the enjoyable collaboration on the generalized VMD approach and many colleagues, including Mary Alberg, Betsy Beise, Stan Brodsky, Gerhard Ecker, Klaus Goeke, Gottfried Holzwarth, Fritz Klein, Mike Musolf, Antoni Szczurek and Wolfram Weise for discussions. This work was supported by the U.S. Department of Energy under Grant No. DE–FG02–93ER–40762 and by the European Community through the HCM programme.

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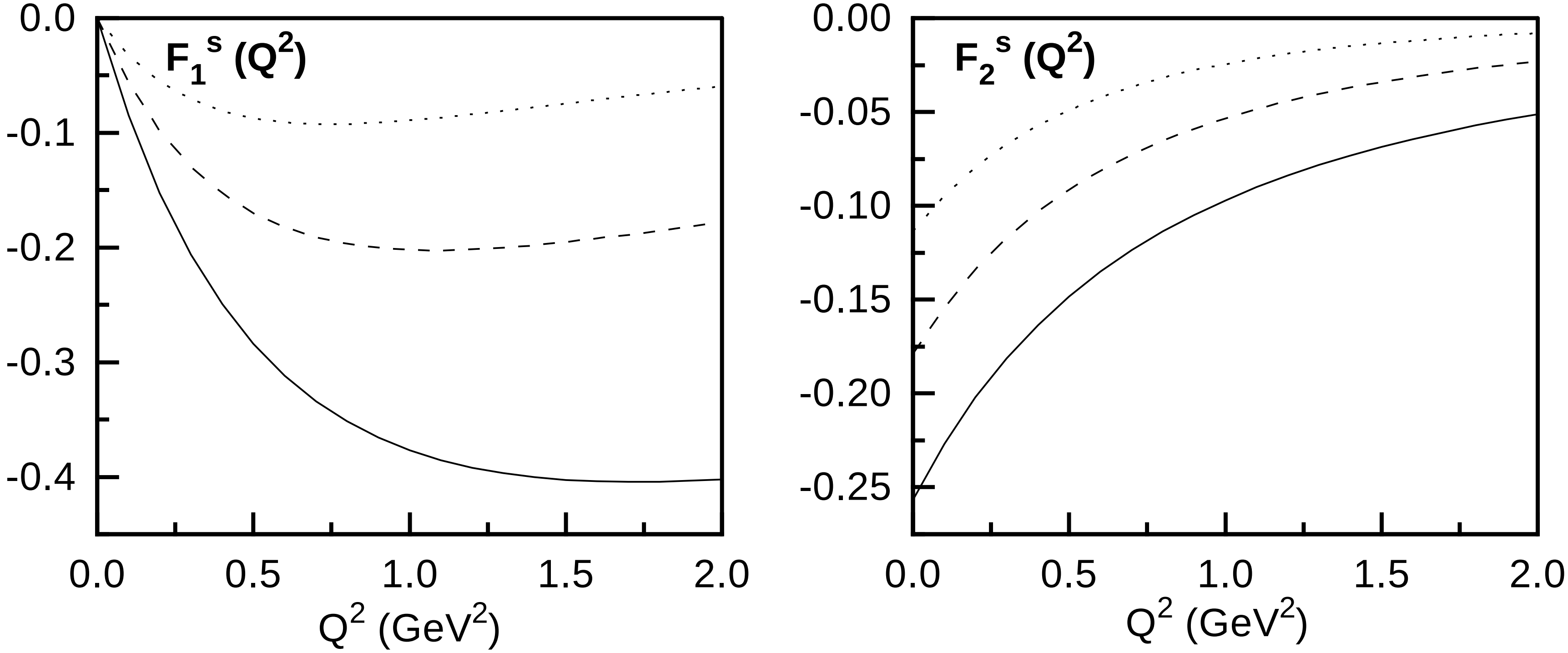
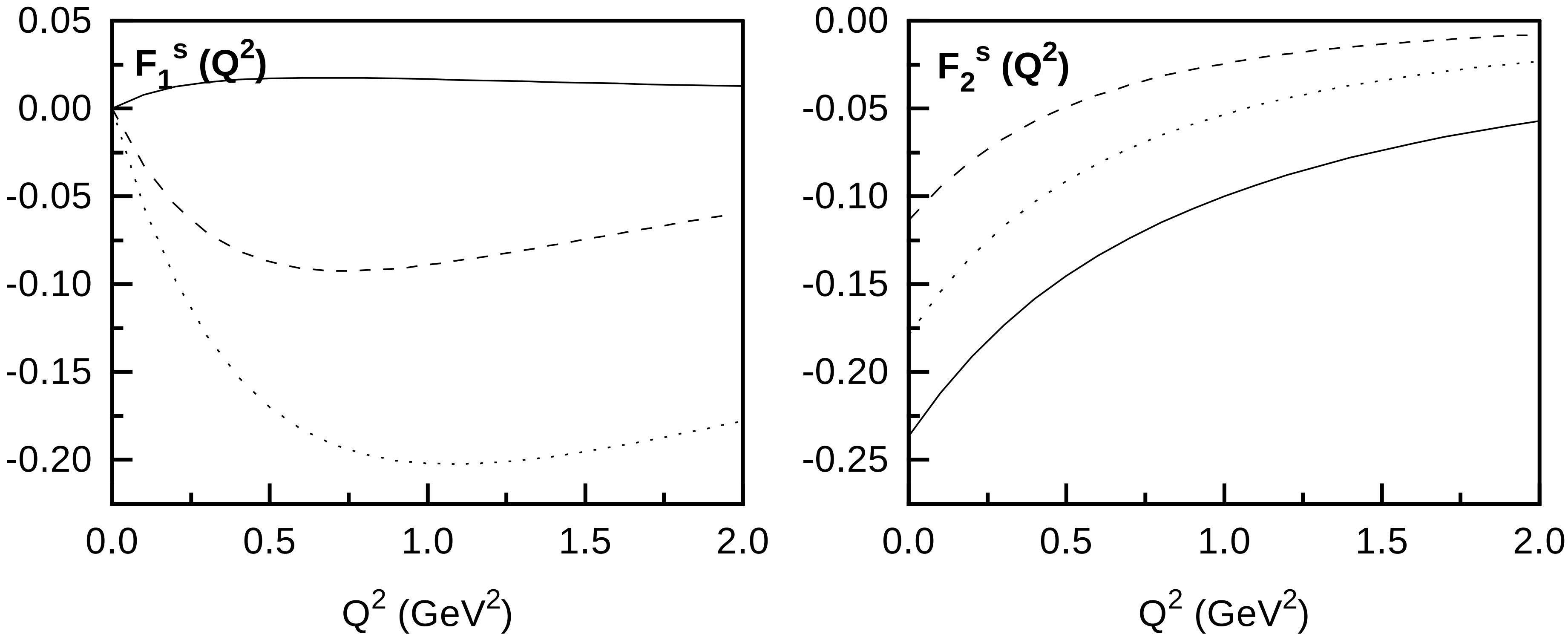
Figure 1: The strange a) Dirac and b) Pauli vector form factors from the 3-pole (continuous line), 4-pole (dashed line) and 5-pole (dotted line) ansätze.

Figure 2: The a) Dirac and b) Pauli form factors from generalized VMD (full line) in comparison with those from the 4-pole (dotted line) and 6-pole (dashed line) ansatz.



$F_1^s\ (Q^2)$
0.0
-0.1
-0.2
-0.3
-0.4
0.0
0.5
1.0
1.5
2.0
Q^2 (GeV^2)
$F_2^s\ (Q^2)$
0.00
-0.05
-0.10
-0.15
-0.20
-0.25
0.0
0.5
1.0
1.5
2.0
Q^2 (GeV^2)



F1s (Q2)
0.05
0.00
-0.05
-0.10
-0.15
-0.20
0.0
0.5
1.0
1.5
2.0
Q2 (GeV2)
F2s (Q2)
0.00
-0.05
-0.10
-0.15
-0.20
-0.25
0.0
0.5
1.0
1.5
2.0
Q2 (GeV2)